

STABLE NON-BPS D-BRANES AND THEIR CLASSICAL DESCRIPTION ¹

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Abstract

We review how to describe the stable non-BPS D-branes of type II string theory from a classical perspective, and discuss the properties of the space-time geometry associated to these configurations. This is relevant in order to see whether and how the gauge/gravity correspondence can be formulated in non-conformal and non-supersymmetric settings.

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1 Introduction

D-branes have played a crucial role in almost all recent developments of string theory. This is mostly due to the fact that they admit a simple and explicit description in terms of open strings with Dirichlet boundary conditions; moreover, being BPS saturated objects, they allow to obtain relevant information about non-perturbative aspects of string theory. In the last couple of years, however, after a remarkable series of papers by A. Sen [1, 2, 3, 4, 5], a lot of attention has been devoted also to the study of non-BPS D-branes (for reviews see Ref.s [6, 7, 8, 9]). There are at least two main motivations which justify this kind of studies. The first one is that non-BPS D-branes are useful in testing various string dualities beyond the BPS level in order to put them on a stronger and more concrete basis. The second motivation, which is even more ambitious, has to do with possible extensions of the Maldacena gauge/gravity correspondence to non-conformal and less supersymmetric or even non-supersymmetric settings. To this purpose, a preliminary but necessary step is represented by the study of the classical geometry generated by the non-BPS branes. In this paper we review the recent results which have been achieved in this direction, mainly referring to Ref. [10].

2 Stable non-BPS D-branes

It is by now well-known that type II string theories in ten dimensions possess non-BPS Dp -branes, with p odd for type IIA and p even for type IIB. These non-BPS D-branes do not preserve any supersymmetry and are unstable due to the existence of open string tachyons on their world-volumes. Indeed, they can decay either into lower dimensional stable objects (*e.g.* BPS D-branes) or into the closed string vacuum depending on which configuration the tachyons acquire. The instability of these non-BPS D-branes clearly puts severe limitations on their use in finding a non-supersymmetric extension of the gauge/gravity correspondence, since a stable geometric background seems to be a crucial and necessary ingredient to this purpose. As a matter of fact, for the non-BPS D-branes of type II in flat space one

cannot disentangle their non-supersymmetric nature from their instability; therefore, despite their intrinsic interest, it is not immediately obvious to say whether or not they can yield a classical geometry and so their use and applications remain doubtful.

Luckily, however, there exist orientifolds and orbifolds of type II theories which admit *stable* non-BPS D-branes! This happens whenever the orientifold/orbifold projection is able to remove any tachyonic state from the open string spectrum on the brane world-volume, so that the resulting configuration is stable even if it is not supersymmetric. For example, in the type I theory, which is an orientifold of the type IIB string with respect to the world-sheet parity, it turns out that the D-instanton and the D-particle are stable non-BPS branes [3, 11]. On the other hand, the existence of such D-branes is required by the type I/SO(32) heterotic string duality, and indeed their properties and interactions are in agreement with this duality [12]. Another well-studied model which possesses stable non-BPS D-branes is the type IIA string on the orbifold T_4/\mathcal{I}_4 , where T_4 is a 4-torus and \mathcal{I}_4 is the parity along its directions. This theory is a singular limit of the type IIA string compactified on a K3 manifold and is related by a non-perturbative duality to the heterotic string compactified on T_4 . From such a duality it is possible to infer that the IIA orbifold model possesses several non-BPS but stable D-branes, like for example the non-BPS D-string [3, 13]. If one performs a T-duality along one of the four compact directions, one obtains the type IIB theory on the orbifold $T_4/(-1)^{F_L}\mathcal{I}_4$, where $(-1)^{F_L}$ is the contribution to the space-time fermion number from the left-moving closed string sector. This is the theory which we consider in the remaining part of this paper.

Let us take for instance a non-BPS Dp -brane of type IIB (*i.e.* p even) whose world-volume is completely transverse with respect to the orbifold directions; then it can be shown that the tachyonic NS ground state of the open string living on the brane world-volume, is globally *odd* under the orbifold parity and is projected out. The lightest states of the would-be tachyon field which can survive the orbifold projection are therefore the windings modes of level one. In fact, these modes are odd under $(-1)^{F_L}$ (because for open string states this operator has the same action as the GSO parity $(-1)^F$), and hence combinations of level-one winding states which are odd under \mathcal{I}_4 are globally even and survive the orbifold projection. More precisely, there are four such states that correspond to the four possible parity-odd combinations along the four compact directions, which we take to be $\{x^6, x^7, x^8, x^9\}$, and that are defined as follows

$$|\chi_a\rangle \equiv |k^\alpha; 1_a\rangle - |k^\alpha; -1_a\rangle \quad (a = 6, \dots, 9) \quad (1)$$

where k^α is the momentum along the flat directions and $\pm 1_a$ represents positive or negative winding of level one along the a -th direction. The mass of these states is given by

$$M_a^2 = \frac{R_a^2}{\alpha'^2} - \frac{1}{2\alpha'} \quad (2)$$

where R_a is the compactification radius of the a -th coordinate. Thus we easily see that for

$$R_a \geq \sqrt{\alpha'/2} \equiv R_{crit} \quad (3)$$

the lightest state of the would-be tachyon field that is selected by the orbifold parity has a non-negative mass squared. In this case, no tachyonic excitation is present on the world-volume and hence the non-BPS Dp -brane is truly stable. A similar analysis holds in various T -dual situations. Clearly, when the T -duality is performed along an orbifolded direction, there is an exchange between windings and momenta, and the above bound gets exchanged to $R_a \leq \sqrt{2\alpha'}$ for the given direction.

The stability of these non-BPS states can be understood also from a closed string point of view. The reason why type II non-BPS D-branes in ten dimensions are unstable is essentially due to the fact that they are neutral objects which are not charged under any R-R field: therefore, their tension is not balanced by any charge “repulsion” which could stabilize them, as it happens for BPS D-branes. Upon orbifold compactification, while these states remain neutral under R-R *untwisted* fields, they actually acquire a charge under *twisted* R-R ones, and when the bound of Eq. (3) is satisfied their charge is sufficient to balance their tension and stabilizes them.

Our next task is to use these non-BPS branes as elementary sources to produce a classical geometry. To this aim, the stability which we have discussed above, is of course a necessary but not sufficient requirement. In fact, one needs also to have big charges, in string units, in order to ensure small curvatures and hence the possibility of neglecting higher derivative terms in the low energy effective theory. In other words one should be able to construct a bound state made of a large number N of microscopic constituents, and this is quite difficult for non-BPS D-branes. However, a remarkable property holds at the border of the stability region defined by Eq.(3): in fact, when the compactification radii are tuned to the critical value R_{crit} , the world-volume states of Eq.(1) become massless and an accidental bose-fermi degeneracy occurs on the non-BPS D-brane [14]. As a consequence, the one-loop open string partition function vanishes. From the closed string point of view, this property is expressed as the vanishing of the tree level exchange amplitude between two boundary states

$$Z = \langle Dp | \mathcal{P} | Dp \rangle = 0 \quad (4)$$

where \mathcal{P} is the closed string propagator and $|Dp\rangle$ is the boundary state representing the Dp -brane as a source of closed string states (for a review on the boundary state formalism and its applications, see Ref.[15]). This means that at critical radii the force between two non-BPS D-branes vanishes, at one loop. In this case, then, the non-BPS D-branes enjoy two properties: they are stable and a no-force condition holds. This clearly opens-up the possibility of constructing macroscopic bound states with these microscopic objects and obtain a classical geometry associated to them. More precisely, it is natural to consider a superposition of a large number of

these stable non-BPS D-branes and simply take the naïve sum of N “single particle” boundary states, that is

$$|Dp, N\rangle \equiv N |Dp\rangle \quad (5)$$

for large N . This is the working hypothesis which we are going to test in the next section.

3 The non-BPS D-particle solution

In this section we consider as a concrete example the non-BPS D-particle of type IIB on $T_4/(-1)^{F_L}\mathcal{I}_4$ and look for the corresponding classical geometry. However, all relevant features of our analysis are to a large extent independent of the specific example considered and are shared by essentially all stable non-BPS D-branes arising in (four dimensional) orbifold compactifications of type II string theories.

The classical description of the non-BPS D-particle can be derived by solving the equations of motion that describe the dynamics of the supergravity fields emitted by the D-particle itself. In this case, such equations are those of the non-chiral $\mathcal{N} = (1, 1)$ supergravity in six dimensions. However, not all fields of this theory are coupled to the D-particle, and to see which ones should be considered, we use the boundary state formalism [15]. In this framework, the non-BPS D-particle is described by a generalized coherent state $|D0\rangle$ which has a NS-NS untwisted component and a R-R twisted one. Its structure is thus of the form

$$|D0\rangle = |D0\rangle_{\text{NS-NS}, U} + |D0\rangle_{\text{R-R}, T_I} \quad , \quad (6)$$

where the index $I = 1, \dots, 16$ in the twisted part indicates on which orbifold plane the D-brane is placed. By computing the overlaps of $|D0\rangle$ with the perturbative closed string states, one gets two essential information: *i*) which are the bulk fields emitted by the brane; *ii*) the asymptotic behavior at large distances of these fields [16]. From this information, the full action describing the dynamics of the fields coupling to the non-BPS D-particle can be easily reconstructed and it turns out to be given by the sum of a bulk action, which is a consistent truncation of the $\mathcal{N} = (1, 1)$ supergravity Lagrangian, and a world-volume action, which acts as a source term.

The explicit form of the coherent state $|D0\rangle$ in Eq.(6) and their overlaps with perturbative closed string states have been studied for example in Ref.s [17, 10]. The result of this analysis is that the six dimensional fields that couple to the non-BPS D-particle are: the graviton $G_{\mu\nu}$, five scalars, *i.e.* φ (related to the ten dimensional dilaton) and η_a ($a = 1, \dots, 4$) (related to the internal components of the ten dimensional metric), and finally one (twisted) vector field A_μ . The action describing the dynamics of these fields at critical compactification radii is

$$\begin{aligned} S = & \frac{1}{2\kappa_{orb}^2} \int d^6x \sqrt{-\det G} \left[\mathcal{R}(G) - \partial_\mu \varphi \partial^\mu \varphi - \partial_\mu \eta_a \partial^\mu \eta_a - \frac{1}{4} e^\varphi F_{\mu\nu} F^{\mu\nu} \right] \\ & - M \int d\tau e^{-\frac{1}{2}\varphi - \frac{1}{2}\sum_a \eta_a} \sqrt{-G_{00}} + M \int A_{(1)} \quad , \end{aligned} \quad (7)$$

where the first line refers to the bulk contribution, the second line to the boundary contribution, and $\kappa_{orb}^2 = 32\pi^3\alpha'^2g^2$, g being the closed string coupling constant. Some comments are in order at this point:

i) The boundary action is not the most general one, since all world-volume fields have been set to zero. However, this is the simplest possible choice which is consistent with our working hypothesis encoded in Eq.(5).

ii) The constant M in front of the DBI term is equal to NM_0 where M_0 is the mass of a single non-BPS D-particle ($M_0 \sim (\sqrt{\alpha'}g)^{-1}$). This is again a consequence of Eq.(5).

iii) Despite the absence of a “mass=charge” relation for a non-BPS brane, the same constant M appears in front of the DBI and WZ terms of Eq.(7). This fact occurs only at the critical radii and can be understood as follows. The relative normalization of the gravitational and gauge terms of the world-volume action can always be adjusted with suitable rescalings of the various fields. In particular one can always make the coefficient of the WZ term become equal to the coefficient of the DBI term by rescaling the R-R potential; but if one does this, in general the corresponding kinetic term in the bulk action acquires a non-canonical normalization. However, at critical radii the rescaling which yields the same number in front of the DBI and WZ terms, is precisely the same which also gives canonically normalized fields in the bulk action. Hence, in this sense one can speak of a “mass=charge” relation even for a non-BPS configuration. This fact can be regarded as the supergravity counterpart of the bose-fermi degeneracy occurring on the brane’s world-volume at critical radii.

We have now all ingredients to see whether a classical description for stable non-BPS D-branes is possible. Although the field equations derived from the action (7) describe a non-BPS configuration, quite surprisingly it is possible to explicitly write the solution in a simple and closed form. In fact, assuming a static and spherically symmetric *Ansatz*, flat asymptotic space geometry, and vanishing asymptotic values for the scalar and gauge fields, the non-BPS D-particle solution turns out to be [10]

$$\varphi = \frac{1}{4} \ln \left[1 + \sin \left(\frac{Q}{r^3} \right) \right] \quad (8)$$

$$\eta_a = \frac{1}{4} \frac{Q}{r^3} \quad (9)$$

$$A_0 = -1 + \frac{\cos \left(\frac{Q}{r^3} \right)}{1 + \sin \left(\frac{Q}{r^3} \right)} \quad (10)$$

$$G_{00} = - \left[1 + \sin \left(\frac{Q}{r^3} \right) \right]^{-3/4} \quad (11)$$

$$G_{ij} = \delta_{ij} \left[1 + \sin \left(\frac{Q}{r^3} \right) \right]^{1/4} \quad (12)$$

where $Q \equiv 2M\kappa^2/3\Omega_4 \sim Ng\alpha'^{3/2}$, and the indices i, j label the transverse directions.

This solution is well defined only for $r > Q^{1/3} \sim (gN)^{1/3}\sqrt{\alpha'}$ and exhibits a naked singularity at

$$r_0 = \left(\frac{2Q}{3\pi}\right)^{1/3} \sim 0.6 Q^{1/3} \quad (13)$$

where the scalar curvature diverges. Within a string theory perspective, naked singularities are allowed if they occur at distances $\leq \sqrt{\alpha'}$ where the supergravity approximation is no longer expected to be valid. From the definition of Q , we see that this is what happens for a single D-particle ($N = 1$): the singularity indeed shows-up at a substringy scale and thus the solution (8)-(12) is acceptable. Moreover, at this scale, the low energy theory becomes effectively ten-dimensional since the compact dimensions are of the same order of r_0 (see Eq.(3)), so one is not only neglecting stringy corrections but the full Kaluza-Klein tower of massive states. The validity of the low-energy solution for a single non-BPS D-particle is consistent with the fact that such state, from a string theory perspective, is stable and expected to exist. However, as we mentioned earlier, a reliable and smooth classical geometry should correspond to big charges, *i.e.* $N \gg 1$. In this case the singularity shows-up in a region where the supergravity approximation is expected to be valid, since now $r_0 \sim Q^{1/3} \sim (gN)^{1/3}\sqrt{\alpha'}$ and $gN \gg 1$, and thus the multi D-particle solution should be rejected. One can say that while its long-distance behavior is consistent and well defined, the supergravity description is probably missing some hidden phenomenon occurring at distances $r \sim r_0$.

There is another feature of our solution that shows that something is missing. In fact, the no-force condition, which holds at one-loop, is broken at higher loops. To understand this point one should keep in mind that Eq.(4) corresponds to a one-loop computation in the open string coupling constant (or, better, in the 'tHooft coupling $\lambda \sim Ng$ which is proportional to Q). The supergravity approximation instead is valid in a very different regime since it is perturbative in α' but *exact* in λ , and it captures the low-energy contribution at *all* loops in the open string coupling. Thus to compare with Eq.(4), we must take the limit $Q \sim \lambda \rightarrow 0$. If we expand our solution (8)-(12) to first order in Q and substitute the resulting expressions into the boundary action, we do find a vanishing force. Indeed, subtracting the vacuum energy, we have

$$\begin{aligned} S_{\text{boundary}} &= -M \int d\tau e^{-\frac{1}{2}\varphi - \frac{1}{2}\sum_a \eta_a} \sqrt{-G_{00}} + M \int d\tau A_0 \\ &\sim -M \int d\tau \frac{Q}{r^3} \left(-\frac{1}{8} - \frac{1}{2} - \frac{3}{8} + 1\right) = 0 \quad , \end{aligned} \quad (14)$$

as expected. A similar calculation however shows that the no-force condition is *not* satisfied at next-to-leading orders, giving evidence that the result of Ref. [14] is spoiled at higher loops. The same kind of conclusion has been derived from a different perspective, namely from a world-volume analysis, in Ref. [18]. As we have mentioned earlier, the world-volume Lagrangian describes the dynamics of a set of fields whose spectrum becomes supersymmetric at the critical radii. Nevertheless,

the Lagrangian itself is not supersymmetric and the interactions break the boson-fermi degeneracy, thus spoiling the one-loop result of Eq.(4). In fact, by taking into account higher loop contributions one can show that non-BPS branes described by a world-volume action like that of Eq.(7) repel each other [19].

4 Discussion and conclusions

We have shown that it is possible to provide a supergravity solution for stable non-BPS D-branes of the type II string on a K3 orbifold in the simplest setting suggested by the one-loop result of Eq.(4) and described by Eq.(5). However, the solution represented by Eqs (8)-(12) has two crucial drawbacks: it is singular, and the no-force condition does not hold when one takes into account higher loop contributions in the open string coupling. These facts indicate that one should go beyond the present analysis, and modify the working hypotheses.

Actually, the two drawbacks are not on the same footing and are not necessarily related to each other. As mentioned above, to have a reliable classical solution within a string context, it is necessary to construct a macroscopic bound state of single microscopic constituents. Since for the non-BPS D-particles the no-force condition does not hold beyond one loop, the appropriate macroscopic bound state must differ from the naïve superposition defined by Eq.(5). To find the correct bound state is therefore the first problem one has to solve. One possibility would be to consider some non-trivial world-volume dynamics, which can radically change the structure of the D-brane and its coupling to the bulk fields. This would lead to an effective action sensibly different from that of Eq.(7), and to new field equations. In Ref. [19], by taking into account string loop corrections to the effective potential of the world-volume theory, it has been found out that for the stable non-BPS D-particle the vacuum at $\chi_a = 0$ (which is the one considered here) is a *local* minimum of the effective potential V while the absolute minimum occurs for (some) $\chi_a \neq 0$, where the non-BPS D-particles attract each other. This fact indicates that within this setting it could be possible to construct a macroscopic bound state corresponding to $N \gg 1$ non-BPS D-branes and study the corresponding classical geometry. Notice that the above result would not only translate into a change of the boundary action but also of the bulk action. In fact, at $\chi_a \neq 0$ the non-BPS states are in the D/\bar{D} fractional-brane phase and therefore couple also to R-R untwisted and NS-NS twisted fields. Therefore, the corresponding consistent truncation of the full $\mathcal{N} = (1, 1)$ supergravity action will be different from the one leading to the action of Eq.(7). Furthermore, one might wonder whether following this strategy, it is possible to automatically solve also the second drawback affecting the solution (8)-(12), namely its singular behavior. This possibility of course deserves and requires further investigation.

In the previous section we have mentioned the fact that the supergravity theory could miss some new physics occurring at distances $r \sim r_0$. A closer look at our singular solution shows that the singularity is actually a *repulson* [20]: the

gravitational force vanishes at some distance $r_e > r_0$ and in the region $r_0 < r < r_e$ the gravitational force is repulsive! It is by now well-understood that this property is shared by many other brane solutions which are dual to non-conformal gauge theories, independently of the amount of preserved supersymmetries [21, 22, 23, 24, 25, 26, 27, 28, 29]. In Ref.[26] a very interesting stringy mechanism has been proposed to excise the singularity and yield a regular solution: independently of the specific details, it turns out that the appropriate source is not point-like in the transverse space, as one could have expected, but rather it is smeared-out on an hypershell called *enhançon* locus. In other words, the constituent branes are forced to cover uniformly an hypersphere rather than pile up in a single point. In the simplest cases the *enhançon* is located precisely at $r = r_e > r_0$ and thus the expanded source surrounds the singularity at $r = r_0$. While the supergravity solution remains unchanged for $r > r_e$, it gets drastically modified in the interior and eventually the singularity is removed. This mechanism has been proved to work in various cases, like for example in the supergravity solutions which are dual to renormalization group flows from $\mathcal{N} = 4$ to $\mathcal{N} = 2$ or $\mathcal{N} = 1$ super Yang-Mills gauge theories [30], or in fractional branes of type II orbifolds [27, 28, 29] (from which the non-BPS D-branes can be obtained as bound states). All these examples are supersymmetric, and it would be very interesting to see whether the *enhançon* mechanism can work also for non-BPS configurations¹. For this, however, it is mandatory to have a better understanding of the low-energy dynamics on the world-volume of non-BPS D-branes, also at loop level. Possibly, the results of [19] could help also in this respect: as already noticed, the true minimum of the stable non-BPS system is claimed to occur at $\chi_a \neq 0$ which corresponds to the D/\bar{D} fractional-brane phase. In this phase the structure of the boundary state is more easily accessible and the relation to the pure supersymmetric fractional brane case is manifest.

When studying any kind of gauge/gravity correspondences, it is necessary to have some control on either side of the duality. This is the reason why one has to understand which are the microscopic objects giving rise to a given supergravity solution. Nevertheless, novel supergravity solutions are also interesting in their own right. Let us then end with a pure supergravity remark. The supergravity field equations one has to solve starting from a bulk action like that of Eq.(7), are second order differential equations and their general solution depends on a certain number of integration constants. These are uniquely fixed by imposing some general physical requirements (*e.g.* asymptotic flat geometry, spherical symmetry, etc.) and by the specific form of the world-volume action, which acts as a source term. Allowing for different types of bound states therefore corresponds to repeat the above analysis and relax those constraints imposed by the boundary action. Under these assumptions, one can see that the corresponding general solution depends on some free parameters related to the mass, the charge and the dilaton and scalar

¹A non-supersymmetric analogue of the *enhançon* mechanism has been recently discussed in [31].

couplings of the given source. Moreover, the periodic functions in Eq.s (8)-(12) get replaced by hyperbolic ones and the general solution reads

$$e^{\eta_a} = \left(\frac{f_-(r)}{f_+(r)} \right)^\delta \quad (15)$$

$$e^{2\varphi} = \left(\frac{\cosh X(r) + \gamma \sinh X(r)}{\cosh \alpha + \gamma \sinh \alpha} \right) \left(\frac{f_-(r)}{f_+(r)} \right)^{\frac{3}{4}\epsilon} \quad (16)$$

$$A_0 = \sqrt{2(\gamma^2 - 1)} \left(\frac{\sinh X(r) (\cosh \alpha + \gamma \sinh \alpha)}{\cosh X(r) + \gamma \sinh X(r)} - \sinh \alpha \right) \quad (17)$$

$$G_{00} = - \left(\frac{\cosh X(r) + \gamma \sinh X(r)}{\cosh \alpha + \gamma \sinh \alpha} \right)^{-\frac{3}{2}} \left(\frac{f_-(r)}{f_+(r)} \right)^{\frac{3}{8}\epsilon} \quad (18)$$

$$G_{ij} = \delta_{ij} \left(\frac{\cosh X(r) + \gamma \sinh X(r)}{\cosh \alpha + \gamma \sinh \alpha} \right)^{\frac{1}{2}} (f_-(r))^{\frac{2}{3} - \frac{1}{8}\epsilon} (f_+(r))^{\frac{2}{3} + \frac{1}{8}\epsilon} \quad (19)$$

where

$$f_\pm(r) = 1 \pm x \frac{Q}{r^3} \quad , \quad X(r) = \alpha + \beta \ln \left(\frac{f_-(r)}{f_+(r)} \right) \quad (20)$$

and $\alpha, \beta, \gamma, \delta, \epsilon$ and x are integration constants which obey

$$\epsilon = \pm \frac{4}{3} \sqrt{4 - 3\beta^2 - 12\delta^2} \quad . \quad (21)$$

If in the general solution (15)-(19) one imposes the behavior dictated by the boundary state (5) or equivalently by the boundary action of Eq.(7), then one obtains the singular solution of Eq.s (8)-(12), but there are other choices of the parameters which instead lead to singularity-free solutions (for details see Ref.[10]). It would be very interesting to investigate further these regular solutions and eventually find which are their microscopic stringy constituents, if any.

The highly non-trivial role that stable non-BPS D-branes can play to obtain a non-conformal and non-supersymmetric extension of the AdS/CFT correspondence makes all this kind of investigations quite challenging.

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